Leptogenesis as the source of gravitino dark matter and density perturbations

Rouzbeh Allahverdi[a] and Manuel Drees[b]

[a] Theory Group, TRIUMF, 4004 Wesbrook Mall, Vancouver, British Columbia, V6T 2A3, Canada.
[b] Physikalisches Institut der Universität Bonn, Nussallee 12, D-53115, Bonn, Germany.

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We investigate the possibility that the entropy producing decay of a right-handed sneutrino condensate can simultaneously be the source of the baryon asymmetry, of gravitino dark matter, and of cosmological density perturbations. For generic values of soft supersymmetry breaking terms in the visible sector of $1 - 10$ TeV, condensate decay can yield the dark matter abundance for gravitinos in the mass range 1 MeV to 1 TeV, provided that the resulting reheat temperature is below 10^6 GeV. The abundance of thermally produced gravitinos before and after sneutrino decay is then negligible. We consider different leptogenesis mechanisms to generate a sufficient asymmetry, and find that low-scale soft leptogenesis works most naturally at such temperatures. The condensate can easily generate sufficient density perturbations if its initial amplitude is $\sim \mathcal{O}(M_{\rm GUT})$, for a Hubble expansion rate during inflation $> 10^9$ GeV. Right-handed sneutrinos may therefore at the same time provide a source for baryogenesis, dark matter and the seed of structure formation.

I. INTRODUCTION

The recent WMAP measurement of the temperature anisotropy of the cosmic microwave background (CMB) has led to the precise determination of many cosmological parameters [1]. Among these, the density of baryons $\Omega_{\rm b}$ and total matter density $\Omega_{\rm m}$ are of particular interest and importance (Ω_X denotes the energy density of species X normalized to the critical energy density of the universe). For the WMAP data only, the best fit parameters are $\Omega_{\rm b} h^2 = 0.024 \pm 0.001$ and $\Omega_{\rm m} h^2 = 0.14 \pm 0.02$, where the present value of the Hubble parameter $h = 0.72 \pm 0.05$ (in units of 100 km sec^{-1} Mpc^{-1}) [1]. The inferred value of $\Omega_{\rm b}$ implies that we live in a baryon asymmetric universe where $\eta_{\rm B} \equiv (n_{\rm B} - n_{\bar{\rm B}})/s \simeq 0.9 \times 10^{-10}$, s being the entropy density of the universe. This is in good agreement with an independent determination from Big Bang nucleosynthesis (BBN) [2]. The determined value of $\Omega_{\rm m}$, on the other hand, implies that most of the matter in universe is non-baryonic and dark. Since a period of inflation [3] washes away any existing baryon and matter density, it seems inevitable that both baryons and dark matter should be generated in the post-inflationary epoch. This has been the focus of intensive research activities in both particle physics and cosmology for the past two decades.

The production of the baryon asymmetry from a baryon symmetric universe requires that three conditions are met: B violation, C and CP violation, and departure from thermal equilibrium [4]. Actually, $B + L$ violating sphaleron transitions, which are active at temperatures 100 GeV $\lesssim T \leq 10^{12}$ GeV [5], imply that any mechanism operating at $T > 100$ GeV must generate a $B - L$ asymmetry. The final baryon asymmetry is then given by $B = a(B - L)$, where $a = 28/79$ in case of the standard model (SM) and $a = 8/23$ for the minimal supersymmetric standard model (MSSM) [6].

Leptogenesis is an attractive mechanism for producing a $B - L$ asymmetry [7]. This scheme postulates the existence of right-handed (RH) neutrinos N_i with a lepton number violating Majorana mass M_N; since the N_i are singlets under the SM gauge group, M_N can be much larger than the electroweak scale. This provides an elegant explanation for the small masses of the light neutrinos via the see-saw mechanism [8]. Moreover, a lepton asymmetry can be generated from the out-of-equilibrium decay of the RH neutrinos, provided CP violating phases exist in the neutrino Yukawa couplings; this lepton asymmetry will be partially converted to a baryon asymmetry via sphaleron effects [7,9,10].

The RH neutrinos can be produced thermally or non-thermally in the early universe. In the thermal scenario [10], the generation of an acceptable lepton asymmetry requires the mass M_1 of the lightest RH neutrino and the temperature of the thermal bath to exceed 10^8 GeV [11–13] (unless RH neutrinos are degenerate [14], or for specific neutrino mass models [15]). However, in many supersymmetric theories this is marginally compatible with the upper bound from thermal gravitino production [16] on the reheat temperature $T_{\rm R} \leq (10^6 - 10^9)$ GeV [17].[*] Alternatively, RH neutrinos can be produced through the perturbative [21] or non-perturbative [22,23] decay of the inflaton; in this case the reheat temperature can be significantly below M_N. Non-thermal leptogenesis can also be achieved without exciting on-shell RH neutrinos [24–26].

In supersymmetric models one also has the RH sneutrinos which can serve as an additional source for leptogene-

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[*] Non-thermal production of gravitinos during preheating [18] does not give rise to severe bounds in realistic models of inflation [19,20].

sis [27]. They are produced along with the neutrinos in a thermal bath or during reheating, and with much higher abundances in preheating [28]. Moreover, there are two unique possibilities for leptogenesis from RH sneutrinos. First, they can acquire a large vacuum expectation value (VEV) if their mass during inflation is less than the Hubble expansion rate at that epoch H_I. This condensate starts oscillating once $H \simeq M_N$, thereby automatically satisfying the out-of-equilibrium condition. The decay of the sneutrino condensate can then yield the desired lepton asymmetry in the same fashion as neutrino decay does [29,30]. The second possibility is to generate a lepton asymmetry in the RH sneutrino sector which will be transferred to the light (s)leptons upon its decay. This asymmetry can be produced via sneutrino couplings to the inflaton [31,32], or from soft supersymmetry breaking effects [33–38].

As already stated, the WMAP measurement also implies the existence of a large amount of non–baryonic dark matter (DM) in the universe. Other evidence for the existence of DM comes from the analysis of galactic rotation curves, from determinations of the masses of clusters of galaxies (i.e. using the X–ray temperature or gravitational lensing techniques), and from attempts to model structure formation in the universe [39]. In supersymmetric models with exact R parity the most natural particle physics DM candidate is the lightest superparticle (LSP). Since superparticles are odd under R parity while "ordinary" particles are even, exact R parity implies that the LSP is stable.

The most widely studied LSP candidate, in particular in the context of dark matter, is the lightest neutralino $\tilde{\chi}_1^0$ [40]. It has several appealing features. Its interactions with ordinary matter, while small, might be sufficient to allow for its detection [40]. Moreover, the $\tilde{\chi}_1^0$ mass, as well as the parameters that determine its couplings, can be measured at colliders [41] (once they reach sufficient energy to produce on-shell superparticles). For some ranges of these parameters $\tilde{\chi}_1^0$ has the right thermal relic density. All these statements also hold in simple, constrained versions of the MSSM, which allow to describe supersymmetry breaking with only a few free parameters [42].

However, neutralinos decouple from the thermal bath at a relatively low temperature $\sim m_{\tilde{\chi}_1^0}/20$. This makes it rather difficult to find a common origin for the creation of the baryon asymmetry and of dark matter. Such a common origin might be hinted at by the observation that Ω_b and Ω_{DM} are of similar order of magnitude. Even ignoring this coincidence, it would certainly be more economical to find a common mechanism for explaining two seemingly unrelated observations.

One such explanation is based on the late decaying Q−ball scenario in models with gravity-mediated supersymmetry breaking [43] (for a similar scenario in gauge-mediated models, see [44]). It is known that in many cases coherent oscillations of flat directions carrying a non-zero baryon number fragment into non-topological solitons, called supersymmetric Q−balls [45]. If the initial VEV of the flat direction is sufficiently large, Q−ball decay (which typically releases three LSPs per unit of baryon number) can occur below the LSP freeze-out temperature [46]. This would give the right DM density for $m_{\tilde{\chi}_1^0} \simeq 2m_p$, which is well below the current experimental lower bound in constrained supersymmetric models [42]. In this scenario overproduction of dark matter can only be avoided if the LSP is dominantly a wino or Higgsino, *and* if the Q−ball decay temperature is tuned to be sufficiently high to allow for significant $\tilde{\chi}_1^0$ annihilation (but below the temperature where $\tilde{\chi}_1^0$ is fully thermalized) [47].

The LSP might also reside in the "hidden" or "secluded" sector, thought to be responsible for the spontaneous breakdown of supersymmetry. Here the most widely studied candidate is the gravitino. Its interactions are determined uniquely by its mass (and the soft breaking parameters in the visible sector). If produced thermally, the gravitino mass density increases linearly with the reheat temperature T_R and inversely with the gravitino mass $m_{3/2}$ [16]. Gravitinos are also produced non-thermally from the decay of the next-to-lightest superparticle (NLSP), typically the lightest neutralino or a scalar lepton. One then has to require that these decays do not violate constraints from BBN [48]. This constraint is easier to satisfy if the NLSP is not a bino. By adding thermal and non–thermal gravitino production, the gravitino can have the required relic density over quite wide ranges of the parameter space [49]. It is even possible to find some combinations of T_R and $m_{3/2}$ that give successful thermal leptogenesis and thermal gravitino dark matter [50]. However, in this scenario the connection between Ω_b and Ω_{DM} is very indirect.

Recently, it has been shown [30,51] that the decay of a RH sneutrino condensate which dominates the energy density of the universe can result in successful leptogenesis as well as dark matter production for gravitino masses in the MeV range. Such small gravitino masses are expected in models with gauge mediated supersymmetry breaking [52], but can also occur in models with gravity mediated supersymmetry breaking for certain (non-canonical) choices of the Kähler metric [53]. In this scenario the late decay of the lightest sneutrino, while generating the lepton asymmetry, reheats the universe and gives rise to thermal production of gravitino dark matter. Obtaining the correct abundances of baryon asymmetry and and dark matter then determine the reheat temperature of the universe $T_R \simeq 10^6$ GeV and the gravitino mass $m_{3/2} \simeq 10$ MeV, respectively [51].

In this note we explore the prospect of scenarios offering an even closer connection between the baryon and gravitino densities, where both are produced *directly* from (s)neutrino decay. The gravitino channel will be kinematically open if supersymmetry breaking results in a mass difference between the RH sneutrino and neutrino

which is $> m_{3/2}$. This mass splitting is essentially determined by the bilinear soft breaking parameter B associated with the (supersymmetric) RH neutrino Majorana mass term. We find that gravitino dark matter with mass $10 \text{ keV} \lesssim m_{3/2} \lesssim 10 \text{ TeV}$ can be produced from sneutrino decay for B in the range $1 - 10$ TeV, compatible with the expected size of visible sector soft breaking terms, provided that the resulting reheat temperature $T_R < 10^6$ GeV. Additional considerations reduce this range to 1 MeV $\lesssim m_{3/2} \lesssim$ 1 TeV. As we will see, among different leptogenesis mechanisms, low-scale soft leptogenesis is preferred to generate the observed asymmetry at such relatively low values of T_R. Sufficient dilution of the gravitinos which are thermally produced at earlier stages typically requires that the initial VEV of the sneutrino condensate be $\simeq 10^{16} - 10^{17}$ GeV. Fluctuations in the sneutrino energy density can then generate cosmological density perturbations via the curvaton mechanism [54,55] if the Hubble expansion rate during inflation is $\sim 10^{11}$ GeV. A RH sneutrino condensate may therefore address three major cosmological issues simultaneously.

The rest of this paper is organized as follows. In Section II we will discuss gravitino production in sparticle decays, specializing to the case for a RH sneutrino. We will turn to leptogenesis from a sneutrino condensate in Section III. In Section IV we will address generation of cosmological density perturbations from a sneutrino condensate. Additional issues such as effects of inflaton decay on the condensate dynamics and vice versa will be considered in Section V. We will conclude by summarizing our results in the closing Section VI.

II. GRAVITINO PRODUCTION FROM SNEUTRINO DECAY

We work in the framework of the MSSM [56] augmented with three RH neutrino multiplets in order to accommodate neutrino masses via the see-saw mechanism [8]. The relevant part of the superpotential is

$$W \supset \frac{1}{2} M_N \mathbf{NN} + \mathbf{h} \mathbf{H}_u \mathbf{NL}, \qquad (1)$$

where $\mathbf{N}$, $\mathbf{H}_u$, and $\mathbf{L}$ are supermultiplets containing the RH neutrinos N and sneutrinos $\tilde{N}$, the Higgs field giving mass to the top quark and its superpartner, and the left–handed (s)lepton doublets, respectively. h are the neutrino Yukawa couplings; for simplicity, family indices on M_N, h, $\mathbf{N}$, and $\mathbf{L}$ are omitted. We work in the basis where the Majorana mass matrix is diagonal.

In addition to the supersymmetry conserving part of the scalar potential for $\tilde{N}$, one also has soft terms from low-energy supersymmetry breaking [56]

$$V_{\text{soft}} \supset m_0^2 |\tilde{N}|^2 + (B M_N \tilde{N}^2 + A h \tilde{N} H_u \tilde{L} + \text{h.c.}), \quad (2)$$

where m_0 and $|B|$ typically are $\mathcal{O}(1 \text{ TeV})$. Supersymmetry breaking by the energy density of the universe also makes contributions $\propto H$ to the soft terms [57]. However, these will be negligible, since $H \ll 1$ TeV during the relevant epoch when the lightest RH (s)neutrino decays, as we will see.

In models of local supersymmetry, the gravitino appears as spin-3/2 partner of the graviton. It acquires a mass $m_{3/2}$ after supersymmetry breaking through the superHiggs mechanism [56]. Gravitino couplings to other particles are suppressed by inverse powers of the reduced Planck mass $M_P \simeq 2.4 \cdot 10^{18}$ GeV, and so is the rate for its decay to or production from matter and gauge fields and their superpartners. Depending on the model, the gravitino mass could be as small as 1 meV or as large as tens of TeV. Not surprisingly, the role the gravitino plays in cosmology depends on the gravitino mass.

If the gravitino is not the LSP, it will decay to particle-sparticle pairs at a rate $\sim m_{3/2}^3/M_P^2$. In the early universe, gravitinos are produced mainly via scatterings of gauge and gaugino quanta in the thermal bath. If $m_{3/2}$ is of the order of the gaugino masses, the helicity $\pm 3/2$ and helicity $\pm 1/2$ states will be produced at approximately the same rate and we will have [16]

$$\frac{n_{3/2}}{s} \simeq 10^{-12} \left(\frac{T_R}{10^9 \text{ GeV}} \right). \qquad (3)$$

For $m_{3/2} \simeq 100 \text{ GeV} - 1 \text{ TeV}$, most gravitinos decay after the onset of BBN and can distort the abundance of light elements synthesized in that epoch. The BBN constraints on the abundance of gravitinos then lead to the bound $T_R \leq (10^6 - 10^9)$ GeV [17]. On the other hand, if $m_{3/2} > 20$ TeV, as in models with anomaly mediated supersymmetry breaking [58], gravitinos decay sufficiently early to evade the BBN constraint. However, LSPs produced in such decays should not overclose the universe. This will result in an upper bound on T_R if gravitino decay occurs below the LSP freeze-out temperature [59]. This bound can be relaxed in those parts of the parameter space which allow a more efficient LSP annihilation and, in consequence, a lower freeze-out temperature [47].

If the gravitino is the LSP, it will be stable, and hence a dark matter candidate, provided that R-parity is unbroken. Then, assuming that gaugino masses are $\gg m_{3/2}$ in this case, the helicity $\pm 1/2$ states are dominantly produced and the fractional energy density of the gravitino will be given by [60,61]

$$\Omega_{3/2} h^2 \simeq \left(\frac{M_{\tilde{g}}}{1 \text{ TeV}} \right)^2 \left(\frac{10 \text{ MeV}}{m_{3/2}} \right) \left(\frac{T_R}{10^6 \text{ GeV}} \right), \quad (4)$$

where $M_{\tilde{g}}$ is the gluino mass. Then, for $M_{\tilde{g}} \sim 500$ GeV, the overclosure bound leads to the constraint $T_R \leq 10^8 m_{3/2}$.

Note that any scenario which attempts to directly link dark matter to leptogenesis can only work if the gravitino is the LSP, or if neutralino LSPs are produced from gravitino decay. This is because leptogenesis can only work at temperature $T \gtrsim 100$ GeV where sphalerons are

still active. At these high temperatures neutralinos are still in thermal equilibrium, so their ultimate relic density will be independent of the number of neutralinos produced from (s)neutrino decay. Here we will consider models with a gravitino LSP where such a link exists. We will treat $m_{3/2}$ and the soft breaking terms in the visible sector, in particular B, as independent free parameters, i.e. we will not assume any specific form for the Kähler metric which determines the relative size of these terms [56].

We want (s)neutrino decays to be the dominant source of relic gravitinos. When the difference between sparticle $\tilde{X}$ and particle X masses $|m_{\tilde{X}} - m_X| \gg m_{3/2}$, which may well be the case if the gravitino is the LSP, helicity $\pm 1/2$ gravitinos are mainly produced in sparticle decays. These states essentially interact like the Goldstinos; the relevant couplings are [60,61]

$$\mathcal{L} \supset \frac{m_{\tilde{X}}^2 - m_X^2}{\sqrt{3} m_{3/2} M_{\mathrm{P}}} \tilde{X}^* \bar{\psi} \left(\frac{1 + \gamma_5}{2} \right) X + \text{h.c.}, \qquad (5)$$

leading to the partial sfermion decay width

$$\Gamma_{\tilde{X} \to X + \psi} \simeq \frac{1}{48\pi} \frac{(m_{\tilde{X}}^2 - m_X^2)^4}{m_{3/2}^2 M_{\mathrm{P}}^2 m_{\tilde{X}}^3}. \qquad (6)$$

Note that a smaller $m_{3/2}$ leads to a more efficient gravitino production.

Here we focus on the production of gravitinos from the decay of the lightest RH (s)neutrino.[*] Since the heavier RH (s)neutrinos essentially play no role in our analysis, we will omit the generation index on the (s)neutrino field. The first two terms in (2) contribute to the mass difference between $\tilde{N}$ and N, but for $M_N \gg 1$ TeV the $B-$term contribution will be dominant. It will result in two sneutrino mass eigenstates $\tilde{N}_1$ and $\tilde{N}_2$, with approximate eigenvalues $M_N - |B|$ and $M_N + |B|$ respectively. Note that B can be made real and positive by a phase transformation of $\tilde{N}$. In that basis, which we will use from now on, $\tilde{N}_2$ and $\tilde{N}_1$ simply are the real and imaginary parts $\tilde{N}_R$ and $\tilde{N}_I$ of $\tilde{N}$. Thus, for $B > m_{3/2}$, the decay channels $\tilde{N}_R \to N + \psi$ and $N \to \tilde{N}_I + \psi$ are kinematically open.

The evolution of $\tilde{N}$ with time has been studied in detail in ref. [35]. Since $\tilde{N}_R$ and $\tilde{N}_I$ evolve independently, the co-moving $\tilde{N}_R$ number density will remain essentially constant for Hubble parameter $H > \Gamma_N$, where $\Gamma_N \simeq |h|^2 M_N/(4\pi)$ is the total $\tilde{N}_R$ decay width. The relative number densities of $\tilde{N}_R$ and $\tilde{N}_I$ are therefore set by the phase of $\tilde{N}$ at the end of inflation. Generically one expects this phase to be $\mathcal{O}(1)$, and hence $\tilde{N}_R$

and $\tilde{N}_I$ to have comparable densities. Here we assume that the sneutrinos dominate the energy density of the universe before they decay. In that case most of today's entropy density originates from sneutrino decays, so that the effective reheat temperature $T_{\mathrm{R}} \simeq \sqrt{M_{\mathrm{P}} \Gamma_N}/2$ (for $g_* \simeq 225$ relativistic degrees of freedom). Moreover, the entropy density just after $\tilde{N}$ decay satisfies $s = 4\rho/(3T_{\mathrm{R}}) \simeq 4 M_N n_{\tilde{N}}/(3T_{\mathrm{R}})$, where ρ is the energy density and $n_{\tilde{N}}$ is the sneutrino number density just before their decay. Then, for $B \gg m_{3/2}$, eq. (6) leads to

$$\left(\frac{n_{3/2}}{s} \right)_{\mathrm{decay}} \simeq 1.6 \cdot 10^{-2} f_R \frac{B^4}{m_{3/2}^2 T_{\mathrm{R}} M_{\mathrm{P}}}, \qquad (7)$$

where we have introduced the quantity $f_R = n_{\tilde{N}_R}/n_{\tilde{N}} < 1$, which we expect to be $\mathcal{O}(1/2)$.

Note that the gravitino abundance from $\tilde{N}_R$ decay is independent of M_N. Since gravitinos are fermions, their occupation number is limited by Pauli blocking to be ≤ 1. In addition, the available phase space in sneutrino decay constrains the physical momentum of produced gravitinos $k_{3/2} \leq B$. This implies an upper limit

$$\left(\frac{n_{3/2}}{s} \right)_{\mathrm{max}} \simeq 3.2 \cdot 10^{-4} \left(\frac{B}{T_{\mathrm{R}}} \right)^3 \qquad (8)$$

on the gravitino abundance. Here we have assumed maximum occupation number throughout the available phase space. This is a valid approximation despite the fact that $k_{3/2}$ is narrowly peaked around B at the time of production. The reason is that the sneutrino decay does not occur instantly. Gravitinos from early sneutrino decays will be redshifted. In the time scale of interest, i.e. $\tilde{N}$ lifetime, $k_{3/2}$ will therefore sweep the phase space due to the Hubble expansion.[†] It can be seen from (7) and (8) that the gravitino abundance will be saturated at its maximum value for

$$m_{3/2} \leq m_{3/2}^{\mathrm{satur}} = 7 \left(\frac{f_R B T_{\mathrm{R}}^2}{M_{\mathrm{P}}} \right)^{1/2}. \qquad (9)$$

On the other hand, for gravitinos to be dark matter, we need

$$\left(\frac{n_{3/2}}{s} \right)_{\mathrm{DM}} \simeq 3 \times 10^{-10} \left(\frac{1 \text{ GeV}}{m_{3/2}} \right). \qquad (10)$$

In order for sphalerons to be active (which is required for leptogenesis) we need $T_{\mathrm{R}} \geq 200$ GeV. On the other hand, eqs.(7) and (10) show that sneutrino decays produce the desired amount of gravitino dark matter if

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[*]Related studies have been performed for gravitino production from the decay of the inflaton [62], the superpartners of stable neutral relics [63], and moduli [64].

[†]In eq.(8) we have assumed that gravitinos do not thermalize. This is true for the combinations of $m_{3/2}$ and T_{R} of relevance to us.

$$m_{3/2} = 0.1 \text{ GeV} \cdot f_R \cdot \frac{200 \text{ GeV}}{T_R} \cdot \left(\frac{B}{1 \text{ TeV}}\right)^4 . \qquad (11)$$

This desired value of the gravitino mass will be above the saturation value (9), if

$$B > 105 \text{ GeV} \cdot \left(\frac{T_R}{200 \text{ GeV}}\right)^{4/7} \cdot \left(\frac{0.5}{f_R}\right)^{1/7} . \qquad (12)$$

For $B = 1$ TeV, this allows reheat temperatures up to 10 TeV. In the framework of leptogenesis, this is still a relatively low temperature. $T_R = 10^6$ GeV can only be tolerated for $B > 14$ TeV. Soft breaking masses of this magnitude are used for first and second generation sfermions in so-called inverted hierarchy models [65]. Much larger values of B do not seem plausible to us. If the constraint (12) is violated, the gravitino relic density, which is now given by the saturation value (8), is too low to account for (all) dark matter.

The lowest possible gravitino mass in this scenario is therefore obtained by setting T_R to its lower bound of 200 GeV and saturating the constraint (12), which gives

$$m_{3/2} > 7 \text{ keV} \cdot \left(\frac{f_R}{0.5}\right)^{3/7} . \qquad (13)$$

It should be noted, however, that gravitino masses below 100 keV lead to over–production of gravitinos from the decay of visible sector superparticles unless the reheat temperature is below the mass scale M_{SUSY} of these superparticles [48]. If we demand that this source of gravitinos contributes at most 10% of the total dark matter density, gravitino masses below 1 MeV would have to be excluded, if $T_R \gtrsim M_{\text{SUSY}}$. Eq. (11) then requires $B > 400$ GeV (for $f_R = 0.5$).

The upper limit of the allowed range of gravitino masses is also reached for $T_R = 200$ GeV, but depends strongly on the largest value of B one is willing to contemplate. For example, taking $B = 20$ TeV, eq.(11) would permit gravitino masses up to 17 TeV. However, since the gravitino is assumed to be the LSP, a gravitino mass well above 1 TeV does not seem plausible.

We shall notice that unlike the more conventional LSP candidates (such as neutralino) there is no prospect for direct detection [40] of gravitino dark matter. Indeed, this holds for all dark matter candidates which have only gravitational interaction with matter [49]. However, indirect detection of DM will in this case be possible through the effects of NLSP decay on BBN, CMB and diffuse photon background [49].

III. LEPTON ASYMMETRY FROM SNEUTRINO DECAY

In the standard (supersymmetric) leptogenesis scenario, the decay of a RH (s)neutrino with mass M_i (we choose $M_1 < M_2 < M_3$) results in a lepton asymmetry per (s)neutrino quanta ϵ_i, given by

$$\epsilon_i = -\frac{1}{8\pi} \frac{1}{[\mathbf{hh}^\dagger]_{ii}} \sum_j \text{Im} \left([\mathbf{hh}^\dagger]_{ij}\right)^2 f\left(\frac{M_j^2}{M_i^2}\right) , \qquad (14)$$

with [66]

$$f(x) = \sqrt{x} \left(\frac{2}{x-1} + \ln\left[\frac{1+x}{x}\right]\right) . \qquad (15)$$

The first and second terms on the right-hand side of eq. (15) correspond to the one-loop self-energy and vertex corrections, respectively. Assuming hierarchical RH (s)neutrinos (so that the asymmetry is produced only from the decay of the lightest RH states), and an $\mathcal{O}(1)$ CP violating phase, it can be shown that [67]

$$|\epsilon_1| \lesssim \frac{3}{8\pi} \frac{M_1(m_3 - m_1)}{\langle H_u \rangle^2}, \qquad (16)$$

where $m_1 < m_2 < m_3$ are the masses of light, mostly left-handed (LH) neutrinos. A hierarchical spectrum for RH (s)neutrinos strongly suggests a hierarchical spectrum of LH neutrino masses, otherwise a big conspiracy would be required to obtain the latter. We will then have

$$|\epsilon_1| \lesssim 2 \cdot 10^{-10} \left(\frac{M_1}{10^6 \text{ GeV}}\right) . \qquad (17)$$

To obtain the last result, we have used $m_3 \simeq 0.05$ eV (as suggested by atmospheric neutrino oscillations) and $\langle H_u \rangle \simeq 170$ GeV. After taking the conversion by sphalerons into account, we arrive at

$$\left(\frac{n_B}{s}\right)_{\text{standard}} \simeq 1.8 \cdot 10^{-10} \left(\frac{T_R}{10^6 \text{ GeV}}\right) . \qquad (18)$$

Obtaining the observed baryon asymmetry from the decay of $\tilde{N}_1$ therefore requires

$$T_R \gtrsim 10^6 \text{ GeV} \qquad (19)$$

even if the (s)neutrinos dominate the energy density of the universe.*

We saw in the previous section that such a high reheat temperature is only marginally compatible with sufficient gravitino production from sneutrino decay. Thus, we should consider other leptogenesis mechanisms which generate a sufficient asymmetry at a lower T_R. One possibility is to have nearly degenerate (s)neutrinos such that $|M_2 - M_1| \ll M_1$. In this case, the $s-$channel pole in

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*In thermal leptogenesis this condition is not satisfied; hence much larger values of M_1 and T_R are needed in that scenario [10,11].

the self-energy correction enhances the generated asymmetry by a factor of $M_1/|\Delta M|$ [14].† Then we can afford to reduce T_R by a factor of $|\Delta M|/M_1$. It is also possible to enhance the produciton asymmetry (17), thus lower T_R, by using a specific neuturino mass model [15]. In this case heavier RH neutrinos N_2 and N_3 make contributions to neutrino masses $\tilde{m}_{2,3}$ which are much larger than the neutrino masses $m_{2,3}$, but cancel each other. Then, for $M_1/M_{2,3} < 100$ and an appropriate pattern of Yukawa couplings, the asymmetry can be sufficiently enhanced to allow successfuil leptogensis for $M_1 \sim \mathcal{O}(\text{TeV})$.

Another, completely different, venue is to consider scenarios which, unlike standard leptogenesis, rely on soft supersymmetry breaking as a source of CP violation. In some of these models, an asymmetry between $\tilde{N}$ and $\tilde{N}^\dagger$ is created by the combined action of neutrino sector $A-$ and $B-$terms [33,34], or from the $B-$term alone via the Affleck-Dine mechanism [35]. This asymmetry is then transferred to the light (s)leptons through $\tilde{N}$ decays as a result of finite temperature supersymmetry breaking. These models can generate a sufficient asymmetry for a rather wide range of M_N, including $M_N \sim 1$ TeV [36], but typically need a suppressed $B-$term $B \ll 1$ TeV. There are also models where the lepton asymmetry is generated in $\tilde{N}$ decay, mainly due to supersymmetry breaking contributions to the vertex correction [37,38].‡ A very interesting observation has been made recently in [38]. There it is shown that the $A-$term and gaugino mass contributions to the vertex correction can lead to successful low-scale leptogenesis, while all soft parameters (including B) assume their natural value $\mathcal{O}(\text{TeV})$. More precisely, if the phase mismatch between A and the electroweak gaugino mass $M_{\widetilde{W}}$ is $\mathcal{O}(1)$, there is a contribution to the asymmetry parameter of order [38]

$$|\epsilon|_{\text{soft}} \sim \alpha_W \frac{|A M_{\widetilde{W}}|}{M_1^2}, \qquad (20)$$

where $\alpha_W \sim 0.03$ is the weak $[SU(2)]$ coupling constant. This implies that

$$\left(\frac{n_B}{s}\right)_{\text{soft}} \lesssim \alpha_W \frac{|A M_{\widetilde{W}}| T_R}{M_N^3}. \qquad (21)$$

Let us choose typical values $|A| \simeq |M_{\widetilde{W}}| \simeq 300$ GeV. Then, even if T_R is near its lower bound of 200 GeV, this mechanism generates a sufficient asymmetry for

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†Of course, one expects zero asymmetry for exactly degenerate (s)neutrinos. The perturbative results in (16) then breaks down, and exact calculations show that no asymmetry will be created [14].

‡All of these models work well with only one generation of (s)neutrinos. This is a marked difference from the standard leptogenesis where mixing among different generations is a necessary ingredient.

$M_N \lesssim 2 \times 10^5$ GeV. A heavier sneutrino will be acceptable for higher reheat temperatures.

Notice that for $M_N \sim \mathcal{O}(\text{TeV})$, soft leptogenesis mechanisms that exploit a $B-$term $\ll 1$ TeV [33–36] can work as well as a source of gravitino dark matter. The reason is that the soft mass m_0 can in this case result in a sufficiently large mass difference between the sneutrino and neutrino. This case can be treated by replacing $B \longrightarrow m_0^2/M_N$ in Eq.(7).

Finally, as noted earlier, we need $T_R > 200$ GeV, so that sphalerons are still active at the time of $\tilde{N}$ decay. Since $\Gamma_N = h^2 M_N/(4\pi) \simeq 4T_R^2/M_P$, this implies for $h^2 \equiv \left[\mathbf{h}\mathbf{h}^\dagger\right]_{11}$:

$$h^2 > 8 \cdot 10^{-18} \frac{10^5 \text{ GeV}}{M_N}. \qquad (22)$$

On the other hand, a sufficient amount of gravitino dark matter could only be generated from sneutrino decay if $T_R \lesssim 10^6$ GeV, which implies the upper bound (note that $T_R < M_N$ for perturbative $\tilde{N}$ decay)

$$h^2 < 2 \cdot 10^{-11} \frac{10^6 \text{ GeV}}{M_N}. \qquad (23)$$

To summarize, soft leptogenesis is the most efficient mechanism for generating a sufficient baryon asymmetry along with gravitino dark matter from the sneutrino decay. It can work for $M_N \lesssim 10^6$ GeV. We will later discuss what this, together with the bounds (22) and (23), implies for the masses of light neutrinos. We now turn to the problem of generating cosmological density perturbations from the sneutrino condensate.

IV. DENSITY PERTURBATIONS FROM SNEUTRINO DECAY

The detection of CMB anisotropies indicates the presence of coherence over superhorizon scales, which is a strong indication of an early inflationary stage [3]. To date measurements are in agreement with the simplest prediction from inflation which is a nearly scale-invariant spectrum of Gaussian and adiabatic primordial perturbations [1]. Traditionally, it has been considered that quantum fluctuations of the inflaton field (which are exponentially stretched during inflation) are responsible for density perturbations [68]. Then, obtaining perturbations of the correct size (about 1 part in 10^5) will constrain parameters of the inflation sector. For example, in the chaotic inflation model with $V(\phi) = m_\phi^2 \phi^2/2$ this leads to $m_\phi \simeq 10^{13}$ GeV.

Models where the inflaton does not generate sufficient perturbations can be rescued, provided that another scalar field is responsible for this. One possibility is that a late-decaying scalar field which acquires inflationary fluctuations dominates the energy density of the universe. Isocurvature fluctuations of this field, called

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curvaton, will in this case be converted to curvature perturbations [54,55].

Another possibility is that the inflaton coupling to matter is controlled by the VEV of some scalar field which acquires fluctuations during inflation. The equation of state then changes at slightly different moments in different parts of the universe at the time of inflaton decay. As a result, this inhomogeneous reheating will also produce curvature perturbations [69–72]. Note that inflationary fluctuations of a scalar field are $\sim H_I$ (recall that H_I is the expansion rate during inflation). Obtaining acceptable perturbations within the curvaton and inhomogeneous reheating mechanisms requires that the scalar field VEV be $\sim 10^5 H_I$.

A RH sneutrino can give rise to density perturbations through either of these mechanisms. The sneutrino can play the role of inflaton in a chaotic inflation model [73,74], in which case $M_N \simeq 10^{13}$ GeV will be needed. Alternatively, for sufficiently small neutrino Yukawas, $\tilde{N}$ can dominate the thermal bath from reheating before it decays, thus playing the role of curvaton [30,75] (for an alternative proposal of a sneutrino as curvaton, see [76]). The inhomogeneous reheating mechanism can also be realized if the inflaton has non-renormalizable couplings to matter fields via $\tilde{N}$, induced by new physics [72,77,78].* An additional advantage when sneutrino is the inflaton or curvaton is that baryon asymmetry of the universe can be directly generated at reheating.

As we saw in the previous sections, simultaneous production of gravitino dark matter and baryon asymmetry from sneutrino decay requires a dominating $\tilde{N}_1$ with mass $M_N \ll 10^{13}$ GeV, which reheats the universe to $T_R \leq 10^6$ GeV. Sneutrino dominance indicates that $\tilde{N}$ is either the inflaton or curvaton.† However, the condition $M_N \ll 10^{13}$ GeV rules out $\tilde{N}_1$ as the inflaton (unless neutrino Yukawas have sufficient fluctuations to allow a successful inhomogeneous reheating).

Now let us find the condition for sneutrino dominance. We start from a $\tilde{N}$ condensate produced during inflation; let $\tilde{N}_0$ be the absolute value of the sneutrino field at the end of inflation. This value, and hence the energy density stored in the sneutrino field, will remain essentially constant ($\rho_{\tilde{N}} \simeq M_N^2 N_0^2$) so long as $H > M_N$. At $H = M_N$ the ratio of the total energy density of the universe ($= 3H^2 M_P^2$) and the energy density in the sneutrino field is therefore $N_0^2/(3M_P^2)$. At that time the sneutrino could therefore only dominate the energy density if $N_0 > M_P$, which is difficult to achieve in the context

of supergravity models. Here we will focus on scenarios with $N_0 < M_P$.

For Hubble parameter $H < M_N$, the energy density of the sneutrinos and their decay products will be redshifted $\propto R^{-3}$ and $\propto R^{-4}$, respectively, where R is the scale factor of the universe. The same is true for the inflatons and their decay products, which, as we just saw, dominate at $H = M_N$. The sneutrino can therefore only become dominant if the inflatons decay before the sneutrinos do, i.e. if the inflaton decay width $\Gamma_\phi > \Gamma_N$. The universe will then go through a radiation dominated epoch, characterized by the temperature $T_I \simeq \sqrt{\Gamma_\phi M_P}/2$ at $H \simeq \Gamma_\phi$ when inflaton decays are completed.

Let us first discuss the case $\Gamma_\phi < M_N$, which seems more plausible for a weakly coupled inflaton. This means that sneutrino oscillations start when the universe is still dominated by inflaton matter. While $M_N > H > \Gamma_\phi$ the amplitude of oscillations will decrease like H; in the radiation dominated epoch, for $H < \Gamma_\phi$, this will change to a behavior $\propto H^{3/4}$, i.e. $\rho_{\tilde{N}} \propto H^{3/2}$. In this radiation dominated epoch the density of the thermal bath (from the inflaton decay products) will decrease like H^2. The two densities will become equal at temperature

$$T_{\rm eq} = T_I \frac{N_0^2}{3M_P^2} . \qquad (24)$$

So far we have ignored $\tilde{N}$ decays. Clearly sneutrino dominance can only occur if $\tilde{N}$ has not yet decayed by the time the inflaton decay products have cooled down to $T = T_{\rm eq}$. This requires $T_{\rm eq} > T_R$; recall that T_R is the temperature after sneutrino decay under the assumptions that sneutrinos do indeed dominate. This condition reads

$$\left(\frac{N_0}{M_P}\right)^2 \geq \frac{3T_R}{T_I}. \qquad (25)$$

A potentially important issue is gravitino production from the thermal bath after reheating. Eq.(4) shows that for values of T_R and $m_{3/2}$ which allow successful leptogenesis and gravitino dark matter from sneutrino decay, thermal production of gravitinos at T_R will be subdominant. It, however, can be significant at the first stage of reheating when $T = T_I$. Gravitino production at this epoch can still be estimated from eq.(4), if we replace T_R by T_I and multiply with the entropy dilution factor from sneutrino decay, which is given by $T_R/T_{\rm eq}$. We require that these thermally produced gravitinos amount to at most 10% of the total dark matter density, i.e. $\Omega_{3/2}^{\rm thermal} h^2 < 0.01$. This leads to the constraint

$$\left(\frac{N_0}{M_P}\right)^2 \geq 0.015 \left(\frac{M_{\tilde{g}}}{1 \text{ TeV}}\right)^2 \frac{10 \text{ MeV}}{m_{3/2}} \frac{T_R}{200 \text{ GeV}} . \qquad (26)$$

The dilution condition (26) is stronger than the one for domination unless there is no gravitino overproduction at T_I. Finally, we can use eq.(11) to express the gravitino mass in terms of B and T_R, to find:

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*Indeed gravitational effects are expected to generate such couplings suppressed by powers of M_P, see Sec. V C.

†A thermally produced $\tilde{N}$ could dominate the universe only if it can be produced through interactions that do not contribute to its decay.

$$\frac{N_0}{M_\mathrm{P}} \geq 0.02 \, \frac{M_{\tilde{g}}}{1 \text{ TeV}} \, \frac{T_\mathrm{R}}{200 \text{ GeV}} \sqrt{\frac{0.5}{f_R}} \left(\frac{1 \text{ TeV}}{B}\right)^2 . \quad (27)$$

Requiring $N_0 < M_\mathrm{P}$ and $M_{\tilde{g}} > 250$ GeV (from collider searches [79]) then implies $B \gtrsim 120$ GeV for our canonical choice $f_R = 0.5$, very similar to the lower bound found in (12). If the RH neutrinos are charged under some extension of the SM gauge groups, as e.g. in an $SO(10)$ model, values of N_0 near the scale of Grand Unification, $M_\mathrm{GUT} \simeq 2 \times 10^{16}$ GeV, seem more natural; this would require $B \gtrsim 2$ TeV. Finally, for $B = 20$ TeV, the constraint (27) allows values of N_0 as small as 10^{14} GeV, if T_R is near 200 GeV. Recall, however, that the combination of large B and small T_R requires quite large gravitino masses, see eq.(11). Since heavy gravitinos have a relatively low thermal relic density, the true bound on N_0 is then often set by the condition (25); in particular, $N_0 = 10^{14}$ GeV is only possible for $T_I \geq 3 \cdot 10^{11}$ GeV.

Generating density perturbations of the correct size via the curvaton mechanism requires [54,55]

$$N_0 \sim 10^5 H_I. \quad (28)$$

For values of N_0 which satisfy (27), this gives $H_I \geq 10^9$ GeV, and we typically have $H_I \simeq 10^{11} - 10^{12}$ GeV. Note that in a chaotic inflation model with such scales inflaton fluctuations do not give rise to sufficient perturbations, and hence another source will indeed be needed.

So far we have assumed that $\tilde{N}$ starts oscillating during the matter dominated epoch. If the first stage of reheating completes when $H > M_N$, the expressions in (24), (25) and (26) will accordingly change. Then, T_I in (24) and (25) should be replaced by the temperature at the onset of sneutrino oscillations $T_\mathrm{osc} \simeq 0.5(M_N M_\mathrm{P})^{1/2}$. Also, the right-hand side of Eq. (26) should be multiplied by a factor of T_I/T_osc. The largest lower bound on N_0 is in this case obtained for an instant inflaton decay, where $T_I \simeq 0.5(H_I M_\mathrm{P})^{1/2}$. After using $H_I \sim 10^{-5} N_0$ and $M_N < 10^5$ GeV so that Eq.(21) generates a sufficient amount of baryon asymmetry, it is seen that $N_0 < M_\mathrm{P} \, (M_\mathrm{GUT})$ requires $B \gtrsim 0.5 \, (2.5)$ TeV. However, smaller values of B are again possible if T_I does not saturate its upper bound.

To summarize, successful sneutrino domination which dilutes thermally produced gravitinos from inflaton decay, along with generating density acceptable perturbations, can be achieved for $N_0 \geq 10^{14}$ GeV and $H_I \geq 10^9$ GeV.

One comment is in order before closing this section. In our scenario sneutrino decay leads to generation of adiabatic baryon and dark matter perturbations. Any amount of baryon asymmetry and dark matter created before sneutrino decay (for example, via thermal leptogenesis and thermal gravitino production) will contribute as an isocurvature component. For a sufficiently late-decaying sneutrino condensate, as considered above, this isocurvature component will be well below the current bounds from WMAP [80]. Future CMB experiments, like PLANCK, will provide a tighter bound on (and might even detect) such a subdominant component.

V. ADDITIONAL CONSIDERATIONS

A. Thermal effects

Since inflaton decay can result in a very high reheat temperature T_I at the end of the first stage of reheating, one has to be careful about thermal effects on the dynamics of the sneutrino condensate. During reheating, the thermal bath consisting of thermalized inflaton decay products has an even higher temperature $T \simeq 0.7(H T_\mathrm{I}^2 M_\mathrm{P})^{1/4} > T_\mathrm{I}$ [81]. During this (inflaton matter dominated) epoch, the temperature therefore scales like $t^{-1/4}$, i.e. $H \propto T^4$, while $H \propto T^2$ after the first stage of reheating is completed.

In our case the only thermal effect one has to worry about is early oscillations of the condensate induced by a large thermal $\tilde{N}$ mass [82], which would dilute the sneutrino density before it decays. If $h N_0 < T$, the Higgs/Higgsino and LH (s)leptons will be in thermal equilibrium before the onset of sneutrino oscillations. This, in turn, induces a thermal mass hT for $\tilde{N}$ which will lead to its early oscillations, if $hT > H$ for some $H > M_N$. Since H grows faster than linearly with T, such thermal oscillations can be avoided if $hT < H$ when $H \simeq M_N$. Thermal effects are most severe if the universe is radiation-dominated at this moment, i.e. if the inflation decay width satisfies $\Gamma_\phi > M_N$, in which case $H(T) \simeq 5T^2/M_\mathrm{P}$. The sufficient condition to prevent early oscillations then reads

$$h^2 \leq 2 \cdot 10^{-13} \frac{M_N}{10^5 \text{ GeV}} , \quad (29)$$

which is somewhat stronger than the constraint (23). However, the limit (29) can be evaded if $h N_0 > T_\mathrm{max}$. Moreover, it becomes weaker by a factor $2\sqrt{M_\mathrm{P} M_N}/T_I$ in the perhaps more realistic case $\Gamma_\phi < M_N$, which implies $T_I < \sqrt{M_\mathrm{P} M_N}/2$.

We note that the upper bound (29) implies that scattering between thermal quanta and the $\tilde{N}$ condensate [35] are not in equilibrium, i.e. the sneutrino field will indeed oscillate coherently until it decays. In any case, the soft leptogenesis mechanism of [38] works just as well for an incoherent ensemble of RH sneutrinos.

Finally, we should mention the possibility of nonperturbative "preheating" decay of the sneutrino condensate [35,83], which can rapidly transfer energy density from the $\tilde{N}$ condensate to a plasma (including $\tilde{N}$ quanta) with energy up to $\sqrt{h N_0 M_N}$, which might exceed M_N significantly [84,85]. It has the undesirable consequence that the sneutrino energy density in this case redshifts faster than that of a coherently oscillating condensate; moreover, since this mechanism prefers bosonic final states, if would suppress gravitino production. As

noted in [35], thermal effects can kinematically block pre-heating, provided that $h N_0 M_N < T^2$ at the onset of $\tilde{N}$ oscillations. Again this constraint is more stringent for $\Gamma_\phi > M_N$, where it implies

$$h < \frac{M_{\mathrm{P}}}{5 N_0} . \tag{30}$$

This bound is much weaker than the one in (23) for all $N_0 \lesssim M_{\mathrm{P}}$.

B. Initial VEV of the sneutrino

As shown in Sec. IV, we need $\tilde{N}$ to have a rather large value $N_0 \gtrsim 10^{14}$ GeV at the end of inflation. In general the evolution of $\tilde{N}$ during inflation is determined by the interplay between quantum fluctuations in deSitter spacetime [3] and the classical equation of motion determined by the sneutrino potential. Here supersymmetry breaking plays an important role, since it generates a contribution $C_I H_I^2 |\tilde{N}|^2$ to the potential [57].

A large value of N_0 is most easily generated if $C_I < 0$. In this case $\tilde{N}$ will quickly settle in the minimum of the potential, which is determined by C_I as well as the terms that stabilize the potential for large $|\tilde{N}|$. For example, if N transforms non-trivially under the GUT gauge group, we expect $N_0 \sim M_{\mathrm{GUT}}$, as mentioned earlier. For $C_I < 0$, $N_0 \sim 10^5 H_I$, as required if $\tilde{N}$ is to serve as curvaton, can be achieved even in "minimal" inflation, which only lasts some 60 e-folds.

On the other hand, $C_I > 1$ would imply $N_0 \lesssim H_I$, which is not sufficient for our purpose. Finally, if $|C_I| \ll 1$, the $\tilde{N}$ mass during inflation will still be given by M_N. For a sufficiently long period of inflation, $N_0 \simeq H_I^2/5 M_N$ will then be obtained [3] at the end of this epoch, regardless of the initial value of $\tilde{N}$. In this case $N_0 \sim 10^5 H_I$ implies $M_N \simeq 2 \cdot 10^{-6} H_I \simeq 2 \cdot 10^{-11} N_0$. Recall from the discussion of Sec. III that soft leptogenesis requires $M_N \lesssim 10^6$ GeV, which implies $N_0 < 2 \cdot 10^{17}$ GeV in this scenario, comfortably in the range of values we found in Sec. IV. However, additional positive terms in the potential could reduce N_0. The value $2 \cdot 10^{17}$ GeV is therefore an upper bound for N_0 if $C_I \geq 0$.

C. Sneutrino VEV and inflaton decay

Sneutrino domination for $N_0 \ll M_{\mathrm{P}}$ requires that the inflatons decay rather early. This corresponds to a high reheat temperature T_I for the first stage of reheating. In particular, $T_I \geq 10^{10}$ GeV, if the inflaton decays before $\tilde{N}$ starts oscillating (assuming that $M_N \geq 1$ TeV). Obtaining such high reheat temperatures requires that some inflaton couplings be sufficiently large. Couplings of the inflaton to other fields are model-dependent and vary from case to case. Recently, it has been noticed

that efficient reheating does not necessarily require large fundamental couplings in the Lagrangian [78]. A scalar condensate with a sufficiently large VEV can result in a large effective coupling of the inflaton, even if fundamental couplings are M_{P} suppressed. Here, we focus on the realization of this scenario, called enhanced reheating, in the presence of the sneutrino condensate. Indeed it will be interesting that $\tilde{N}$ can set the stage for its domination by facilitating inflaton decay.

To elucidate, we consider a simple superpotential term

$$W \supset \lambda \frac{\Phi}{M_{\mathrm{P}}} \mathbf{N H}_u \mathbf{L}, \tag{31}$$

where Φ is the inflaton superfield. Regardless of any discrete or continuous global symmetry, gravitational effects are expected to generate such a term in the superpotential. If Φ remains a gauge singlet up to M_{P}, $\lambda \sim \mathcal{O}(1)$ is expected. On the other hand, if the inflaton is non-singlet under some extended gauge group which is broken at a scale M_{new}, λ will be suppressed by powers of $M_{\mathrm{new}}/M_{\mathrm{P}}$.

The rate for inflaton decay from (31) reads

$$\Gamma_\phi \simeq \frac{\lambda^2}{4\pi} \left(\frac{N_0}{M_{\mathrm{P}}} \right)^2 m_\phi. \tag{32}$$

For chaotic inflation, $m_\phi = H_I$, and the condition for acceptable perturbations via curvaton mechanism $N_0 \sim 10^5 H_I$ leads to

$$T_{\mathrm{I}} \simeq 4 \times 10^{-4} \lambda \frac{N_0^{3/2}}{\sqrt{M_{\mathrm{P}}}}. \tag{33}$$

If the inflaton is an absolute singlet, $\lambda \simeq 1$. For $N_0 \simeq M_{\mathrm{GUT}} \simeq 2 \cdot 10^{16}$ GeV, we will then have $T_{\mathrm{I}} \sim 7 \cdot 10^{11}$ GeV. Therefore, a sneutrino VEV which is typically preferred by our scenario, can naturally lead to a fast inflaton decay as needed in the curvaton mechanism.

D. Non-dominating sneutrino

So far our results have been derived for the case when the sneutrino condensate dominates the universe before decaying. This is necessary if the sneutrino is to serve as curvaton. In this section we investigate if RH sneutrinos that do *not* dominate the energy density of the universe can simultaneously produce leptogenesis and gravitino dark matter. Since there will be only one stage of reheating, it will be more appropriate to replace T_{R} and T_{I} with the sneutrino decay temperature* T_{d} and the universe reheat temperature T_{R}, respectively. Note that T_{R}

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*Note that this is now simply the temperature at which the sneutrino decays, i.e. the temperature of the thermal plasma at $H = \Gamma_N$. By assumption this temperature is not changed significantly by these decays. We still assume that the sneutrino decays after the inflaton does.

should now be sufficiently low so that thermal gravitino production (4) is subdominant.

The sneutrino energy density at $H \simeq \Gamma_N$ will now be a fraction $r < 1$ of the energy density of the thermal bath from inflaton decay. For $\Gamma_\phi < M_N$[†] this fraction can be computed from

$$r = \frac{T_\mathrm{R} N_0^2}{3 T_\mathrm{d} M_\mathrm{P}^2}. \tag{34}$$

The number density of gravitinos produced in sneutrino decay is reduced from (7) to (we take $f_R = 0.5$ from now on)

$$\left(\frac{n_{3/2}}{s}\right)_\mathrm{decay} \simeq 8 \cdot 10^{-3} r \frac{B^4}{m_{3/2}^2 T_\mathrm{d} M_\mathrm{P}}. \tag{35}$$

The upper bound on the gravitino number density from phase space is still given by Eq. (8), with $T_\mathrm{R} \to T_\mathrm{d}$. The gravitino density from sneutrino decay will saturate this bound for gravitino mass

$$m_{3/2}^\mathrm{satur} = 5 \left(\frac{r B T_\mathrm{d}^2}{M_\mathrm{P}}\right)^{1/2}. \tag{36}$$

Eq. (35) gives the correct dark matter density for

$$m_{3/2} = 56\,\mathrm{MeV}\, r \left(\frac{B}{1\,\mathrm{TeV}}\right)^4 \frac{200\,\mathrm{GeV}}{T_\mathrm{d}}. \tag{37}$$

This will be above the saturation value (36) if

$$r \geq 1.4 \cdot 10^{-7} \left(\frac{1\,\mathrm{TeV}}{B}\right)^7 \left(\frac{T_\mathrm{d}}{200\,\mathrm{GeV}}\right)^4. \tag{38}$$

The lowest possible value of the gravitino mass compatible with gravitino dark matter from sneutrino decay is obtained when the bound (38) is saturated:

$$m_{3/2} \geq 8\,\mathrm{eV} \left(\frac{1\,\mathrm{TeV}}{B}\right)^3 \left(\frac{T_\mathrm{d}}{200\,\mathrm{GeV}}\right)^3. \tag{39}$$

The constraint (38) would allow values of r below 10^{-14} for $B > 10$ TeV, if $m_{3/2}$ is well below 1 eV. However, as already noted at the end of Sec. II, gravitinos of such low mass would be over-produced from the decays of visible sector sparticles [48]. Requiring $m_{3/2} > 1$ MeV in order to keep this contribution subdominant even if sparticles are light, and again taking $B \lesssim 20$ TeV from naturalness argument, we see from (37) that $r \gtrsim 10^{-7}$ is required.

We now turn to sneutrinos as source of the baryon asymmetry. We saw in Sec. III that standard high scale

leptogenesis, where CP violation comes from the neutrino Yukawa couplings, is only marginally compatible with gravitino DM from sneutrino decay even if sneutrinos dominate the universe, unless two (or more) RH sneutrinos are closely degenerate, or for a specific neutrino mass model. Here we therefore focus on soft leptogenesis. Eq. (21) is modified to

$$\left(\frac{n_\mathrm{B}}{s}\right)_\mathrm{soft} \lesssim \alpha_W r \frac{|A M_{\widetilde{W}}| T_\mathrm{d}}{M_N^3}. \tag{40}$$

Note that this expression is only valid for $M_N^2 > |A M_{\widetilde{W}}|$. Since $\alpha_W \simeq 0.03$, leptogenesis therefore also requires $r \gtrsim 10^{-7}$. This bound can be saturated only for values of M_N not much in excess of 1 TeV. In that case the term $\propto m_0^2$ in the scalar potential (2) can also contribute significantly to the sneutrino-neutrino mass splitting. The desired gravitino density can then be obtained for smaller values of B.

Finally, the curvaton mechanism will not work if $\tilde{N}$ does not dominate. The sneutrino condensate can nevertheless play a crucial role in the generation of density perturbations via inhomogeneous reheating proceeding through the coupling in (31). It is interesting that for a smaller N_0, usually needed to have $r < 1$, this coupling also results in a lower T_R required to avoid thermal overproduction of gravitinos. Of course, we still need $N_0 \simeq 10^5 H_I$ for this mechanism to work. Scenarios with $r \ll 1$, and correspondingly small values of N_0, would therefore require a relatively low Hubble parameter during inflation. For example, taking $T_\mathrm{R} \sim 10^3 T_\mathrm{d}$, eq.(34) shows that $H_I \sim 10^8$ GeV would be required if r is near its lower limit of 10^{-7}.

VI. SUMMARY AND CONCLUSIONS

We have shown that an SU(2) singlet "right-handed" sneutrino field $\tilde{N}$, which has originally been introduced in the context of supersymmetric see-saw models of neutrino masses, can be the simultaneous source of the baryon asymmetry (via leptogenesis), of nonthermal dark matter (via its decay into gravitinos), and of density perturbations (via either the curvaton or inhomogeneous reheating mechanism). This can most easily be achieved if the sneutrinos dominated the total energy density of the universe for some period of time after inflation. Dark matter production from this source requires a relatively low reheat temperature after sneutrino decay ($T_\mathrm{R} \lesssim 10^6$ GeV). This is only marginally compatible with standard leptogenesis, where the necessary CP violation comes from the neutrino Yukawa couplings, but comfortably fits in the range of temperatures where "soft leptogenesis" can work, where the source of CP violation is in the soft breaking terms. On the other hand, our mechanism can work for a wide range of gravitino masses (1 MeV $\lesssim m_{3/2} \lesssim$ 1 TeV), and only imposes a mild lower bound on the soft breaking B parameter associated with

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[†]Recall that early inflaton decays more easily lead to sneutrino domination; hence we consider relatively late inflaton decays here.

the large Majorana mass driving the see-saw mechanism, $B \gtrsim 400$ GeV.

If the sneutrino is not the curvaton, its energy density need not dominate. In fact, it can still explain both the baryon asymmetry and gravitino dark matter if its energy density only amounted to 10^{-7} of the total. In this case we need a lower sneutrino decay temperature, larger B, and/or smaller $m_{3/2}$, compared to the case where the sneutrino dominates.

One perhaps not so attractive aspect of our model (also existing in the model of Refs. [30,51]) is that the Yukawa couplings of $\tilde{N}$ should be small, see (23). Such a small value for h implies that the lightest LH neutrino ν_1 should be extremely light; for example, taking typical values $T_R = 1$ TeV, $M_N = 100$ TeV we find $m_{\nu_1} \sim 10^{-7}$ eV. Theoretically, it has been suggested that a small Yukawa coupling (as well as $M_N \ll M_{GUT}$), can be explained in the context of orbifold GUT models [86]. Experimentally, an extremely light (even massless) ν_1 is allowed and has testable consequences for neutrinoless double beta decay experiments [87]. Therefore a tiny Yukawa coupling, despite being not attractive, can be accommodated by theory and experiment.

The sneutrino energy density can easily be large enough for the sneutrino to serve as curvaton or to trigger inhomogeneous reheating if its Hubble induced soft SUSY breaking mass during inflation was negative or small. In particular, we find acceptable scenarios where the sneutrino field after inflation is of the order of the scale of Grand Unification, as expected e.g. in $SO(10)$ models. If we only want the sneutrinos to produce the baryon asymmetry and gravitino dark matter, sufficiently many $\tilde{N}$ quanta could also have been produced from inflaton decay. However, the sneutrino abundance is model dependent in this case and can vary considerably.

It should be admitted that we have not explained why the baryon and dark matter densities are of comparable magnitude. Even though they have a common origin, their numerical values depend on different combinations of parameters. In particular, the dark matter density scales like the fourth power of B and is inversely proportional to the gravitino mass. These parameters do not affect the baryon density at all. Conversely, the dark matter density does not depend on the CP violating phases that crucially enter the expression of the baryon asymmetry. However, on close inspection this seems to be true for all models that have been proposed so far. We find it encouraging that for the soft leptogenesis scenario the main remaining freedom comes from soft breaking parameters. Within our framework the baryon and dark matter densities can therefore be considered as cosmological constraints on the mechanism of supersymmetry breaking. Even in the absence of an explicit model that gives the required relations between soft breaking parameters we find it encouraging that we can solve three of the most important problems in current cosmology using a single field, which has the additional advantage of being well motivated from particle physics.

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